

D-term Enhancement in Spin-1 Top Partner Model

Haiying Cai

Department of Physics, Peking University, China.

Supersymmetric models with extended electroweak gauge groups have the potential to enhance the Higgs quartic interaction through nondecoupling D-terms. We consider the D-term enhancement effect in a vector top partner model, where the quadratic divergence to the Higgs mass from the virtual top quark is canceled by its corresponding spin-1 superpartners. We are going to show that the model can predict a Higgs mass beyond the LEP bound, and is consistent with the precision electroweak constraints.

I. INTRODUCTION

In the Supersymmetric theory, since the quadratic divergence associated with the Higgs mass-squared from the SM fields will be canceled by their superpartners, soft SUSY breaking terms only induce logarithmic corrections and the scalar field mass is stabilized to be around the soft SUSY breaking scale m_s . In order to let the SUSY theory to be natural and therefore reduce fine-tunings, the soft SUSY breaking scale m_s is supposed to be in the hundred GeV range. In the minimal supersymmetric standard model (MSSM), the physical Higgs mass is related to the mass of Z gauge bosons times a factor of $\cos(2\beta)$ at the tree level, where β is determined by the ratio of the two Higgs fields vacuum expectation values (VEVs) (v_u/v_d). However, the LEP direct search excluded the existence of a Higgs bosons below 114.4 GeV at 95% C.L. For the Higgs to go beyond the LEP bound, large radiative contribution to the quartic interaction term from the top quark sector is necessary, which in turn demands the top squark to have a mass of the TeV order. The tension between the electroweak scale and the new physics emerging scale, which is referred as the little hierarchy problem, encourages people to explore new possibilities to avoid the dilemma. There are many attempts to achieve a Higgs mass much heavier than the Z gauge bosons in the supersymmetric theory. One straightforward way is to enhance the quartic interaction term at the tree level, and generally additional interaction structure is required. In the NMSSM model [1], one extra $SU(2)_L$ singlet superfield N is added, which couples with the two Higgs fields through a supersymmetric Yukawa interaction $\lambda N H_u H_d$. A large λ is preferred to generate a large quartic term but the requirement that the Yukawa interaction is perturbative till the unification scale puts an upper bound on the Higgs mass. An alternative method to raise the Higgs mass without inducing fine tunings is to consider a fat Higgs scenario originated from a strong interaction sector. In the fat higgs scenario, the singlet chiral field N and the two Higgs fields H_u and H_d are composite meson fields interacting via a naturally large Yukawa coupling. The original fat Higgs model has a dynamically generated superpotential $\lambda N(H_u H_d - v^2)$ with the similar matter content as the NMSSM in the low energy scale [2]. This type of theory is further extended by Refs. [3] and [4]. In the New Fat Higgs model [3], only the singlet chiral field N is composite while the two Higgs fields are still kept elementary.

Supersymmetric models with enlarged gauge groups under which the Higgs bosons are charged may raise the Higgs mass through the nondecoupling D-term effects [5] [6]. In the low energy scale, the enlarged gauge groups need to be broken into the Standard Model gauge group by the VEVs of some extra scalar fields. If the gauge symmetry is broken in a SUSY conserving limit, the D-term effects of these extra scalar fields would decouple and we could recover the standard MSSM D-term potential for the Higgs fields. In order to retain the D-term effects from those extra fields till the electroweak scale, SUSY breaking effects need to be included in the mechanism responsible for the gauge symmetry breaking. When the SUSY breaking scale is much larger than the gauge symmetry breaking scale, the effective D-term for the Higgs fields in the electroweak scale can be enhanced.

In this Letter I consider the possibility to increase the Higgs quartic interaction terms in a spin-1 top partner model [7]. In this model, the superpartners of the left-handed top quark are spin-1 vector bosons, while the superpartner of the right-handed top quark is still a scalar. This scenario is realized by extending the gauge group and assembling the left-handed top quark into a vector supermultiplet. The extended gauge group serves to provide the source of nondecoupled D-term effects. Extra chiral fields are necessary to be added to trigger the gauge symmetry breaking since we hope to achieve a D-term flat minimum. In the following of this letter, I will specify the superpotential responsible for gauge symmetry breaking and supersymmetry breaking. The exact mass spectrum for scalar states in the link fields after the symmetry breaking will be calculated. I am going to verify the D-term enhancement effects in the Higgs bosons sector and explore the bound for the mass of the Higgs boson in this model after considering relevant electroweak constraints.

	$SU(3)$	$SU(2)$	$U(1)_H$	$U(1)_V$	$SU(5)$	$H + V + aT_{24}$
H	1	1	$\frac{1}{2}$	$\frac{1}{10}$	\square	$(\frac{2}{3}, \frac{1}{2})$
\bar{H}	1	1	$-\frac{1}{2}$	$-\frac{1}{10}$	$\bar{\square}$	$(-\frac{2}{3}, -\frac{1}{2})$
Φ_3	$\bar{\square}$	1	$-\frac{1}{6}$	$\frac{1}{10}$	\square	$(0, -\frac{1}{6})$
Φ_2	1	$\bar{\square}$	0	$\frac{1}{10}$	\square	$(\frac{1}{6}, 0)$
$\bar{\Phi}_3$	\square	1	$\frac{1}{6}$	$-\frac{1}{10}$	$\bar{\square}$	$(0, \frac{1}{6})$
$\bar{\Phi}_2$	1	\square	0	$-\frac{1}{10}$	$\bar{\square}$	$(-\frac{1}{6}, 0)$

TABLE I: The gauge structure of the two Higgs fields and the four link fields .

II. D-TERM ENHANCEMENT AND HIGGS MASS

We first briefly review the structure of the spin-1 top partner model. For details of the realization, one can refer to the previous paper [7]. The model is based on the gauge group $SU(3) \times SU(2) \times U(1)_H \times U(1)_V \times SU(5)$, which can be better illustrated in a supersymmetric two sites moose diagram (See Fig [1]). One copy of three generations of leptons, quarks plus their superpartners are put in the first moose site with a gauge group of $SU(3) \times SU(2) \times U(1)_H$. These chiral superfields transform exactly the same as in the MSSM. While two higgs superfields H and \bar{H} need to be put in a second moose site which has a $SU(5) \times U(1)_H \times U(1)_V$ gauge group. The gauge coupling of $U(1)_H$ can be set to be very small. Four vector-like link fields $\Phi_3, \bar{\Phi}_3, \Phi_2, \bar{\Phi}_2$ are responsible to communicate between the fields located in the two isolated moose sites. When the link fields gain nonzero VEVs, they break the original product gauge group $SU(3) \times SU(2) \times U(1)_H \times SU(5) \times U(1)_V$ down to the diagonal MSSM gauge group $SU(3)_C \times SU(2)_W \times U(1)_Y$. The gauge transformation property of the Higgs fields and the four link fields are given in Table I.

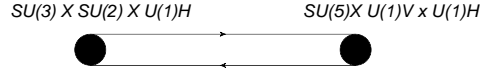


FIG. 1: two sites illustration for spin-1 top partner model

The chiral superfields $\bar{H} = (\bar{T}, H_1)$ and $H = (\bar{T}^c, H_2)$ transform as fundamental and antifundamental under the $SU(5)$ gauge group respectively. H_1 and H_2 are the traditional up-type and down-type Higgs superfields in the MSSM. Since the triplet companies of H_1 and H_2 transform as $(\bar{3}, 1)$ and $(3, 1)$ under the diagonal $SU(3)_C \times SU(2)_W$ subgroup, they are treated as one couple of vector-like right-handed heavy top quarks in our model. The link fields can mix \bar{T} with the right-handed field \bar{u}_3 living in the first moose site, so that one linear combination still has zero mass and thus is the physical right-handed top quark. Both of the $U(1)_H$ and $U(1)_V$ charges are necessary for arranging a correct hypercharge and a correct baryon number for the Higgs bosons and top quarks.

For the vector supermultiplet in the $SU(5)$ gauge group, the extra gauge bosons X, Y transforming as $(3, 2)$ under the $SU(3)_C \times SU(2)_W$ are identified as the spin-1 top partner in this model. They are lifted to be heavy after the gauge symmetry breaking in the same way as in the GUT model. We also need to identify the field contents in the four link fields. Under the diagonal $SU(3)_C \times SU(2)_W$, $\Phi_3 (\bar{\Phi}_3)$ will split into one complex singlet $\Phi_{3S} (\bar{\Phi}_{3S})$, one complex octet $\Phi_{3O} (\bar{\Phi}_{3O})$, plus one component field $\Phi_{3t} (\bar{\Phi}_{3t})$ with the same quantum number as $\bar{t}_L (t_L)$ in the first moose site. Similar decomposition applies to the fields of Φ_2 and $\bar{\Phi}_2$. They will split into one complex singlet $\Phi_{2S} (\bar{\Phi}_{2S})$, one complex triplet $\Phi_{2T} (\bar{\Phi}_{2T})$, as well as one component field $\Phi_{2t} (\bar{\Phi}_{2t})$ which has the same quantum number as $t_L (\bar{t}_L)$. Φ_{2t} and $\bar{\Phi}_{3t}$ mix with the $(3, 2)$ sector gaugino λ_{32} through gauge interaction and they mix with the left handed field $Q_3 = (t_L, b_L)$ in the first moose site through Yukawa interaction. As long as the μ -terms for $\Phi_{2,3}$ and $\bar{\Phi}_{2,3}$ are large enough compared with their VEVs, the dominant component of the physical left handed top quark will be the gaugino λ_{32} and their superpartner are spin-1 gauge bosons.

We now write down the superpotential relevant for calculations. In order to get quartic terms for the link fields, we may add two singlets chiral superfields $S_{1,2}$ to interact with $\Phi_{2,3}$ and $\bar{\Phi}_{2,3}$. One adjoint chiral superfield A_1 charged under the $SU(2)$ gauge group and another adjoint chiral superfield A_2 charged under the $SU(3)$ gauge group are also added to ensure that there are no light modes after the gauge symmetry breaking.

$$\begin{aligned}
W_{susy} = & y_1 Q_3 \Phi_3 \bar{\Phi}_2 + \mu_2 \Phi_2 \bar{\Phi}_2 + \mu_3 \Phi_3 \bar{\Phi}_3 + \lambda_S S_1 (\Phi_2 \bar{\Phi}_2 - w_2^2) + \lambda_S S_2 (\Phi_3 \bar{\Phi}_3 - w_3^2) \\
& + \lambda_A \bar{\Phi}_2 A_1^a \frac{\sigma^a}{2} \Phi_2 + \lambda_A \bar{\Phi}_3 A_2^m G^m \Phi_3.
\end{aligned} \tag{1}$$

$\sigma^a/2, (a = 1, 2, 3)$ are the generators for the $SU(2)$ gauge group and $G^m, (m = 1, \dots, 8)$ are the generators for the $SU(3)$ gauge group. The two λ_S singlet interacting terms will force VEVs for $\Phi_{2,3}$ and $\bar{\Phi}_{2,3}$. The first Yukawa interaction term is not relevant for the gauge symmetry breaking but it will align the VEVs of $\Phi_{2,3}$ and $\bar{\Phi}_{2,3}$ in the singlet component field direction, i.e.

$$\begin{aligned} \langle \Phi_3 \rangle &= \begin{pmatrix} f_3 & 0 & 0 & 0 & 0 \\ 0 & f_3 & 0 & 0 & 0 \\ 0 & 0 & f_3 & 0 & 0 \end{pmatrix}, \quad \langle \bar{\Phi}_3 \rangle^T = \begin{pmatrix} \bar{f}_3 & 0 & 0 & 0 & 0 \\ 0 & \bar{f}_3 & 0 & 0 & 0 \\ 0 & 0 & \bar{f}_3 & 0 & 0 \end{pmatrix}, \\ \langle \Phi_2 \rangle &= \begin{pmatrix} 0 & 0 & 0 & f_2 & 0 \\ 0 & 0 & 0 & 0 & f_2 \end{pmatrix}, \quad \langle \bar{\Phi}_2 \rangle^T = \begin{pmatrix} 0 & 0 & 0 & \bar{f}_2 & 0 \\ 0 & 0 & 0 & 0 & \bar{f}_2 \end{pmatrix}. \end{aligned} \quad (2)$$

We can check from Table I that these singlets' VEVs do not violate the $H + V + aT_{24}$ charge. The VEVs break the large gauge group down into the MSSM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ and their gauge couplings are given by:

$$\frac{1}{g_{2,3}^2} = \frac{1}{\hat{g}_{2,3}^2} + \frac{1}{\hat{g}_5^2}, \quad \frac{1}{g_Y^2} = \frac{1}{\hat{g}_{1H}^2} + \frac{1}{\hat{g}_{1V}^2} + \frac{1}{15\hat{g}_5^2}, \quad (3)$$

where \hat{g}_i and \hat{g}_5 are the gauge couplings of the original $SU(3)$, $SU(2)$, $U(1)_H$, $U(1)_V$ and $SU(5)$ gauge groups respectively. For simplicity, we further assume $f_2 = \bar{f}_2$ and $f_3 = \bar{f}_3$, therefore this is a D-term flat minimum and it will not induce mass terms for the Higgs fields. The singlet field $S_{1,2}$ and adjoint fields $A_{1,2}$ will not gain VEVs in this scenario. These terms give the following scalar potential for the four link fields $\Phi_{2,3}$ and $\bar{\Phi}_{2,3}$:

$$\begin{aligned} V_\Phi &= y_1^2 |\phi_3 \bar{\phi}_2|^2 + \lambda_S^2 |\text{Tr } \phi_2 \bar{\phi}_2 - w_2^2|^2 + \lambda_S^2 |\text{Tr } \phi_3 \bar{\phi}_3 - w_3^2|^2 \\ &+ \mu_2^2 (\text{Tr } \phi_2^\dagger \phi_2 + \text{Tr } \bar{\phi}_2 \bar{\phi}_2^\dagger) + \mu_3^2 (\text{Tr } \phi_3^\dagger \phi_3 + \text{Tr } \bar{\phi}_3 \bar{\phi}_3^\dagger) \\ &+ \lambda_A^2 |\text{Tr } \bar{\phi}_2 (\sigma^a/2) \phi_2|^2 + \lambda_A^2 |\text{Tr } \bar{\phi}_3 G^m \phi_3|^2 \end{aligned} \quad (4)$$

The minimum of this simple potential determines the VEVs,

$$f_2^2 = \bar{f}_2^2 = \frac{\lambda_s^2 w_2^2 - \mu_2^2}{2\lambda_s^2}, \quad (5)$$

$$f_3^2 = \bar{f}_3^2 = \frac{\lambda_s^2 w_3^2 - \mu_3^2}{3\lambda_s^2}. \quad (6)$$

Substituting Eq. [5] and Eq. [6] back into the scalar potential Eq.[4], we can see that Supersymmetry is spontaneously broken in this setup via the ORaifeartaigh mechanism with the simultaneous presence of the supersymmetric μ terms and the λ_S interaction terms. An easy way to verify this statement is that, only if $\mu_2 = \mu_3 = 0$, the superpotential could have zero vacuum energy when link fields develop nonzero VEVs. Desired values for the two VEVs f_2 and f_3 can be achieved by tuning the three free parameters: $\mu_{2,3}$, $w_{2,3}$ and λ_S .

As we expect no light modes after the gauge symmetry breaking, we first examine the mass spectrum in the link fields after they gain VEVs. Due to the traceless properties of the $SU(2)$ and $SU(3)$ gauge generators, the two λ_A terms will not change the VEVs, but they will give masses to the two linear copies of real triplet fields i.e. $\psi_{2T,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2T} + \bar{\phi}_{2T}^\dagger)$ and $\psi_{2T,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{2T} - \bar{\phi}_{2T}^\dagger)$, as well as the two linear copies of real octet fields i.e. $\psi_{3O,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3O} + \bar{\phi}_{3O}^\dagger)$ and $\psi_{3O,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{3O} - \bar{\phi}_{3O}^\dagger)$, leaving other states untouched. The two λ_S quartic terms and the two $\mu_{2,3}^2$ mass terms can give masses to two specific linear copies of real triplet fields i.e. $\psi_{2T,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2T} - \bar{\phi}_{2T}^\dagger)$ and $\psi_{2T,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{2T} - \bar{\phi}_{2T}^\dagger)$, plus two specific linear copies of real octet i.e. $\psi_{3O,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3O} - \bar{\phi}_{3O}^\dagger)$ and $\psi_{3O,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{3O} - \bar{\phi}_{3O}^\dagger)$. They also give masses for six real singlet states: $\psi_{2S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2S} - \bar{\phi}_{2S}^\dagger)$, $\psi_{2S,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2S} + \bar{\phi}_{2S}^\dagger)$, $\psi_{2S,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{2S} - \bar{\phi}_{2S}^\dagger)$; $\psi_{3S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3S} - \bar{\phi}_{3S}^\dagger)$, $\psi_{3S,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3S} + \bar{\phi}_{3S}^\dagger)$, $\psi_{3S,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{3S} - \bar{\phi}_{3S}^\dagger)$. Ignoring some singlets mixing, we list the mass spectrum for all the singlets, triplets and octets in Table II. As shown in that table, we have one copy of real triplet field, one copy of real octet field and two copies of real singlet fields left massless, which are the Goldstone bosons eaten by the heavy W' , G' gauge bosons and two heavy $U(1)$ gauge bosons B' and B'' respectively.

$\frac{1}{\sqrt{2}}\text{Re}(\phi_{2T} - \bar{\phi}_{2T}^\dagger) :$	$2\mu_2^2 + 2(\hat{g}_2^2 + \hat{g}_5^2)f_2^2$	$\frac{1}{\sqrt{2}}\text{Re}(\phi_{2T} + \bar{\phi}_{2T}^\dagger) :$	$\lambda_A^2 f_2^2$
$\frac{1}{\sqrt{2}}\text{Im}(\phi_{2T} - \bar{\phi}_{2T}^\dagger) :$	$(2\mu_2^2 + \lambda_A^2 f_2^2)$	$\frac{1}{\sqrt{2}}\text{Im}(\phi_{2T} + \bar{\phi}_{2T}^\dagger) :$	0
$\frac{1}{\sqrt{2}}\text{Re}(\phi_{2S} - \bar{\phi}_{2S}^\dagger) :$	$2\mu_2^2$	$\frac{1}{\sqrt{2}}\text{Re}(\phi_{2S} + \bar{\phi}_{2S}^\dagger) :$	$(2\lambda_S^2 w_2^2 - 2\mu_2^2)$
$\frac{1}{\sqrt{2}}\text{Im}(\phi_{2S} - \bar{\phi}_{2S}^\dagger) :$	$2\lambda_S^2 w_2^2$	$\frac{1}{\sqrt{2}}\text{Im}(\phi_{2S} + \bar{\phi}_{2S}^\dagger) :$	0
$\frac{1}{\sqrt{2}}\text{Re}(\phi_{3O} - \bar{\phi}_{3O}^\dagger) :$	$2\mu_3^2 + 2(\hat{g}_3^2 + \hat{g}_5^2)f_3^2$	$\frac{1}{\sqrt{2}}\text{Re}(\phi_{3O} + \bar{\phi}_{3O}^\dagger) :$	$\lambda_A^2 f_3^2$
$\frac{1}{\sqrt{2}}\text{Im}(\phi_{3O} - \bar{\phi}_{3O}^\dagger) :$	$(2\mu_3^2 + \lambda_A^2 f_3^2)$	$\frac{1}{\sqrt{2}}\text{Im}(\phi_{3O} + \bar{\phi}_{3O}^\dagger) :$	0
$\frac{1}{\sqrt{2}}\text{Re}(\phi_{3S} - \bar{\phi}_{3S}^\dagger) :$	$2\mu_3^2$	$\frac{1}{\sqrt{2}}\text{Re}(\phi_{3S} + \bar{\phi}_{3S}^\dagger) :$	$(2\lambda_S^2 w_3^2 - 2\mu_3^2)$
$\frac{1}{\sqrt{2}}\text{Im}(\phi_{3S} - \bar{\phi}_{3S}^\dagger) :$	$2\lambda_S^2 w_3^2$	$\frac{1}{\sqrt{2}}\text{Im}(\phi_{3S} + \bar{\phi}_{3S}^\dagger) :$	0

TABLE II: mass spectrum for singlets (ignore their mixing) , triplets and octets in the link fields $\phi_{2,3}$ and $\bar{\phi}_{2,3}$.

The two real singlets fields $\psi_{2S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2S} - \bar{\phi}_{2S}^\dagger)$ and $\psi_{3S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3S} - \bar{\phi}_{3S}^\dagger)$ mix in a more complicated way due to $U(1)$ D-terms proportional to \hat{g}_S^2 , \hat{g}_{1H}^2 and \hat{g}_{1V}^2 . Their mass eigenstates are determined by diagonalizing the full mass terms for $\psi_{2S,1}$ and $\psi_{3S,1}$ as described in Eq. [7] :

$$\frac{2\hat{g}_5^2}{15} \left(3f_2\psi_{2S,1} - \sqrt{6}f_3\psi_{3S,1} \right)^2 + \frac{\hat{g}_{1H}^2}{3} f_3^2 \psi_{3S,1}^2 + \frac{\hat{g}_{1V}^2}{50} \left(2f_2\psi_{2S,1} + \sqrt{6}f_3\psi_{3S,1} \right)^2 + 2\mu_2^2 \psi_{2S,1}^2 + 2\mu_3^2 \psi_{3S,1}^2 \quad (7)$$

The component fields in the off-diagonal direction also need to be dealt with in a straightforward way. For convenience, we conduct an eigenstate basis rotation and redefine those stop-like states in the four link fields as follows:

$$\psi_{2t,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2t} - \bar{\phi}_{2t}^\dagger), \quad \psi_{2t,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2t} + \bar{\phi}_{2t}^\dagger), \quad \psi_{2t,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{2t} - \bar{\phi}_{2t}^\dagger), \quad \psi_{2t,4} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{2t} + \bar{\phi}_{2t}^\dagger), \quad (8)$$

$$\psi_{3t,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3t} - \bar{\phi}_{3t}^\dagger), \quad \psi_{3t,2} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3t} + \bar{\phi}_{3t}^\dagger), \quad \psi_{3t,3} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{3t} - \bar{\phi}_{3t}^\dagger), \quad \psi_{3t,4} = \frac{1}{\sqrt{2}} \text{Im}(\phi_{3t} + \bar{\phi}_{3t}^\dagger). \quad (9)$$

In terms of the new eigenstate basis of $\psi_{2t,i}, i = 1, 2, 3, 4$, and $\psi_{3t,i}, i = 1, 2, 3, 4$, the mass terms from both the gauge interactions and superpotentials are:

$$\hat{g}_5^2 (f_2\psi_{2t,1} + f_3\psi_{3t,1}^T)^2 + 2\mu_2^3 \psi_{2t,1}^2 + 2\mu_3^3 (\psi_{3t,1}^T)^2 + \frac{y_1^2}{2} (f_2(\psi_{3t,1}^T + \psi_{3t,2}^T) - f_3(\psi_{2t,1} - \psi_{2t,2}))^2, \quad (10)$$

$$\hat{g}_5^2 (f_2\psi_{2t,3} - f_3\psi_{3t,3}^T)^2 + 2\mu_2^3 \psi_{2t,3}^2 + 2\mu_3^3 (\psi_{3t,3}^T)^2 + \frac{y_1^2}{2} (f_2(\psi_{3t,3}^T + \psi_{3t,4}^T) + f_3(\psi_{2t,3} - \psi_{2t,4}))^2. \quad (11)$$

The origins of each term in the above two equations are: the first one is from the Super-Higgs mechanism, the second one and third one are from the $\mu_{2,3}$ mass terms as well as the λ_S interaction terms, while the last one comes from the $y_1 Q_3 \Phi_2 \bar{\Phi}_3$ Yukawa interaction. Similar to the singlets case, the exact mass eigenstates are determined by diagonalizing these two equations. Examining Eq.[10] and Eq.[11], it is easy to find out that only two linear combinations of stop-like states are still massless and they should be identified as the Goldstone bosons for the heavy X, Y gauge bosons.

$$\pi_t = \frac{1}{\sqrt{f_2^2 + f_3^2}} (f_2\psi_{2t,2} - f_3\psi_{3t,2}^T) \quad (12)$$

$$\eta_t = \frac{1}{\sqrt{f_2^2 + f_3^2}} (f_2\psi_{2t,4} + f_3\psi_{3t,4}^T) \quad (13)$$

With the mass spectrum for all the scalars fields in $\phi_{2,3}$ and $\bar{\phi}_{2,3}$, we proceed to discuss the effective D-term in this model. In a supersymmetric gauge theory, when a large gauge group breaks down into the MSSM gauge group, the vector supermultiplet corresponding to the unbroken generators will inherit the MSSM gauge interactions and remain massless after the gauge symmetry breaking. For the vector supermultiplet corresponding to the broken generators, they could achieve masses after eating a copy of chiral supermultiplet through the Super-Higgs mechanism. In the supersymmetric limit, all component fields $(A_\mu, \lambda_1, \lambda_2, \Sigma)$ in a heavy vector supermultiplet have degenerate masses. And after integrating them out, the D-term effects from those heavy states will decouple. In order to retain the D-term effects from those heavy states till the low energy scale, a SUSY breaking mass term need to be added to the real scalar component field Σ , i.e. the lowest component field in the heavy vector supermultiplet, which will recouple the D-term effects from the broken gauge generators back into the effective Lagrangian. In the low energy scale, since the Higgs bosons are charged under the diagonal $SU(2)_W$ and $U(1)_Y$ gauge group, there should be two sources of D-term enhancement in this model: One is from an extra $SU(2)$ embedded in the $SU(5)$ gauge group and the other is from two extra $U(1)$ s. For the heavy $SU(2)$ vector supermultiplet, its corresponding scalar component field is a triplet state: $\psi_{2T,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2T} - \bar{\phi}_{2T}^\dagger)$. While for two extra $U(1)$ s, their scalar component fields are two heavy singlets: $\psi_{2S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{2S} - \bar{\phi}_{2S}^\dagger)$ and $\psi_{3S,1} = \frac{1}{\sqrt{2}} \text{Re}(\phi_{3S} - \bar{\phi}_{3S}^\dagger)$. After integrating out those heavy fields, an effective D-term is obtained for the Higgs bosons:

$$\frac{g_2^2}{2} \Delta_2 \left(H_2^\dagger \frac{\sigma^a}{2} H_2 - H_1^\dagger \frac{\sigma^a}{2} H_1 \right)^2 + \frac{g_Y^2}{2} \Delta_Y \left(\frac{1}{2} H_2^\dagger H_2 - \frac{1}{2} H_1^\dagger H_1 \right)^2, \quad (14)$$

where g_2 is the gauge coupling for the SM gauge group $SU(2)_W$ and g_Y is the gauge coupling for the SM hypercharge gauge group $U(1)_Y$, whose values are determined by Eq.[3]. The D-term effects of these heavy scalar fields are nondecoupling due to spontaneous SUSY breaking effects in our scenario, and their effects can be summarized into two parameters Δ_2 and Δ_Y :

$$\Delta_2 = \left(1 + \frac{m_{2T}^2}{f_2^2} \frac{1}{\hat{g}_2^2} \right) / \left(1 + \frac{m_{2T}^2}{f_2^2} \frac{1}{(\hat{g}_5^2 + \hat{g}_2^2)} \right), \quad (15)$$

$$\begin{aligned} \Delta_Y = & \left(1 + \left(\frac{4}{5\hat{g}_5^2} + \frac{1}{2\hat{g}_{1H}^2} + \frac{1}{2\hat{g}_{1V}^2} \right) \frac{m_{2S}^2}{f_2^2} + \left(\frac{1}{5\hat{g}_5^2} + \frac{3}{\hat{g}_{1V}^2} \right) \frac{m_{3S}^2}{f_3^2} \right. \\ & + \frac{15\hat{g}_5^2 + 25\hat{g}_{1H}^2 + \hat{g}_{1V}^2}{10\hat{g}_5^2\hat{g}_{1H}^2\hat{g}_{1V}^2} \frac{m_{2S}^2 m_{3S}^2}{f_2^2 f_3^2} \Bigg) / \left(1 + \frac{60\hat{g}_5^2 + 25\hat{g}_{1H}^2 + 9\hat{g}_{1V}^2}{2(\hat{g}_{1H}^2\hat{g}_{1V}^2 + 15\hat{g}_5^2(\hat{g}_{1H}^2 + \hat{g}_{1V}^2))} \frac{m_{2S}^2}{f_2^2} \right. \\ & \left. + \frac{3(15\hat{g}_5^2 + \hat{g}_{1V}^2)}{(\hat{g}_{1H}^2\hat{g}_{1V}^2 + 15\hat{g}_5^2(\hat{g}_{1H}^2 + \hat{g}_{1V}^2))} \frac{m_{3S}^2}{f_3^2} + \frac{75}{2(\hat{g}_{1H}^2\hat{g}_{1V}^2 + 15\hat{g}_5^2(\hat{g}_{1H}^2 + \hat{g}_{1V}^2))} \frac{m_{2S}^2 m_{3S}^2}{f_2^2 f_3^2} \right). \end{aligned} \quad (16)$$

m_{2T} , m_{2S} and m_{3S} are the respective F-term masses for the heavy scalar fields $\psi_{2T,1}$, $\psi_{2S,1}$ and $\psi_{3S,1}$ induced by spontaneous SUSY breaking, whose values can be read from the first term in each column of Table II:

$$m_{2T}^2 = 2\mu_2^2, \quad m_{2S}^2 = 2\mu_2^2, \quad m_{3S}^2 = 2\mu_3^2 \quad (17)$$

In the supersymmetric limit i.e. $\mu_2 = \mu_3 = 0$, we can find that $\Delta_2 = 1$ and $\Delta_Y = 1$, that is the D-term effects from those heavy fields are decoupled. However in this model, since we prefer to stay in the region of $\hat{g}_5 f_3 \ll \mu_3$ and $\hat{g}_5 f_2 \ll \mu_2$, we expect that there are notable enhancements for both of the $SU(2)_W$ and $U(1)_Y$ D-terms. Under the limit of $m_{2T} \gg f_2$, $m_{2S} \gg f_2$, and $m_{3S} \gg f_3$, Δ_2 and Δ_Y are simply determined by those gauge coupling constants:

$$\Delta_2 \approx \frac{\hat{g}_2^2 + \hat{g}_5^2}{\hat{g}_2^2 \hat{g}_5^2} \hat{g}_5^2 = \frac{\hat{g}_5^2}{\hat{g}_2^2}, \quad \Delta_Y \approx \frac{15\hat{g}_5^2 + 25\hat{g}_{1H}^2 + \hat{g}_{1V}^2}{25\hat{g}_Y^2} \quad (18)$$

We can see that in the large SUSY breaking limit, the effective D-term for the Higgs bosons is only proportional to three gauge coupling constants \hat{g}_5^2 , \hat{g}_{1H}^2 and \hat{g}_{1V}^2 , which is exactly the same as in the original unbroken gauge theory. The Higgs bosons can gain notable mass through the D-term enhancement effects as long as we choose the gauge couplings under which our Higgs bosons are charged to be larger than the gauge couplings in the other moose site. An simple example is choosing $\hat{g}_2 = 0.78$, $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$ and $\hat{g}_{1V} = 1.5$, we will obtain $\Delta_2 \sim 3.36$ and $\Delta_Y \sim 8.21$. With an $O(1) \tan(\beta) \sim 2.0$, at the tree level the Higgs mass squared $M_h^2 \leq \frac{1}{4}(g_2^2 \Delta_2 + g_Y^2 \Delta_Y) v^2 \cos^2(2\beta)$ can be

naturally raised to be around $(115 \text{ GeV})^2$. Since the radiation corrections from the top quark and its superpartners will contribute to the running of the Higgs quartic term :

$$\delta\lambda = \frac{3\hat{g}_5^4}{8\pi^2} \ln\left(\frac{m_{\bar{Q}}m_{\bar{t}_R}}{m_t^2}\right), \quad (19)$$

at the loop level the Higgs mass is further enhanced such that

$$m_h^2 = \frac{1}{2} \left(m_A^2 + M_G^2 - \sqrt{(m_A^2 + M_G^2)^2 - 4m_A^2 M_G^2 \cos^2 2\beta} \right) + 2\delta\lambda v^2 \sin(\beta)^2, \quad (20)$$

with $M_G^2 = \frac{1}{4} (g_2^2 \Delta_2 + g_Y^2 \Delta_Y) v^2$.

The precision of Higgs mass prediction depends on the mass of vector top partner and the mass of the right handed stop used to calculate the radiation correction. The vector top partner gains its mass through the link fields' VEVs, i.e. $m_{\bar{Q}} = \hat{g}_5^2(f_2^2 + f_3^2)$. While the right handed stop acquires its mass through the higgs μ_H term $\mu_H H \bar{H}$ as well as from the soft SUSY breaking scalar mass term. If the vector top partner mass is $m_{\bar{Q}} \simeq 1.9 \text{ TeV}$ and the mass of right handed stop is $m_{\bar{t}_R} \simeq 450 \text{ GeV}$, and let the mass parameter M_A to be 800 GeV , we can get a heavy Higgs bosons $m_h \simeq 195 \text{ GeV}$.

III. ELECTROWEAK CONSTRAINTS: S, T, U AND $Z \rightarrow b\bar{b}$

Those gauge couplings should be chosen so that they could reproduce the SM gauge couplings at the EW scale after running by the renormalization group. In the following, we are going to take some specific sets of gauge couplings when we proceed with the electroweak analysis, so that the electroweak measurements are adopted to constrain the link field VEVs, f_2 and f_3 , whose values in turn determine the masses of the vector top partner and Higgs boson in this model. The most stringent constraints come from the mixing of W, B and W', B', B'' due to the Higgs VEVs because the two Higgs doublets and the light fermions transform under different gauge groups. Oblique corrections to the Standard Model are contained in the vacuum polarizations of gauge bosons, which are parameterized by S, T and U . The vacuum polarizations of a gauge boson can be expanded around the zero momentum.

$$\Pi_{a,a'}(p^2) = \Pi_{a,a'}(0) + p^2 \Pi'_{a,a'}(0) + \dots, \quad (21)$$

and the S, T and U parameters are defined in the following way:

$$S = 16\pi \cdot (\Pi'_{33}(0) - \Pi'_{3Q}(0)), \quad T = \frac{4\pi}{s^2 c^2 M_Z^2} (\Pi_{11}(0) - \Pi_{33}(0)), \quad U = 16\pi \cdot (\Pi'_{11}(0) - \Pi'_{33}(0)). \quad (22)$$

with $s = \sin(\theta)$ and $c = \cos(\theta)$ and θ is Weinberg angle. The definition of S, T, U subtracts out the predicted SM contribution with fixed top quark mass and Higgs bosons mass so that they encode only new physics contributions. The contribution to the S and U parameters from the gauge bosons mixing is very less, which are at the order of v^4/f_3^4 or v^4/f_2^4 . The analytic expressions for S and U are simply given by:

$$\begin{aligned} \Delta S = & -\frac{\pi}{4} \left(\frac{(25\hat{g}_1^2 + 60\hat{g}_5^2 + 9\hat{g}_{1V}^2)(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(-15\hat{g}_5^2 + 4\hat{g}_{1V}^2))}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_2^4} + \frac{\hat{g}_5^4}{(\hat{g}_2^2 + \hat{g}_5^2)^3} \frac{v^4}{f_2^4} \right) \\ & - \frac{9\pi\hat{g}_{1H}^2(150\hat{g}_5^4 - 5\hat{g}_5^2\hat{g}_{1V}^2 - \hat{g}_{1V}^4)(-15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 - 4\hat{g}_{1V}^2))}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_2^2 f_3^2} \\ & - \frac{9\pi\hat{g}_{1H}^4(15\hat{g}_5^2 + \hat{g}_{1V}^2)^3}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_3^4}, \end{aligned} \quad (23)$$

$$\begin{aligned} \Delta U = & \frac{\pi}{8} \frac{(25\hat{g}_1^2 + 60\hat{g}_5^2 + 9\hat{g}_{1V}^2)(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(-15\hat{g}_5^2 + 4\hat{g}_{1V}^2))}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_2^4} \\ & + \frac{9\pi\hat{g}_{1H}^2(150\hat{g}_5^4 - 5\hat{g}_5^2\hat{g}_{1V}^2 - \hat{g}_{1V}^4)(-15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 - 4\hat{g}_{1V}^2))}{2(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_2^2 f_3^2} \\ & + \frac{9\pi\hat{g}_{1H}^4(15\hat{g}_5^2 + \hat{g}_{1V}^2)^3}{2(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^3} \frac{v^4}{f_3^4}. \end{aligned} \quad (24)$$

The experimental constraints for S and U are:

$$S = 0.04 \pm 0.10, \quad U = 0.08 \pm 0.11. \quad (25)$$

If we assume $0.4\text{TeV} < f_3 \ll f_2$, and with a small \hat{g}_{1H} but a large \hat{g}_{1V} , the S and U parameters do not put any constraint to our parameter space. The situation is different for the other oblique parameter. The gauge bosons mixing can give sizable contribution to the T parameter:

$$\Delta T = \frac{1}{\alpha} \left(\frac{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(-15\hat{g}_5^2 + 4\hat{g}_{1V}^2))^2}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{16f_2^2} + \frac{3\hat{g}_{1H}^4(15\hat{g}_5^2 + \hat{g}_{1V}^2)^2}{(15\hat{g}_5^2\hat{g}_{1V}^2 + \hat{g}_{1H}^2(15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{8f_3^2} \right). \quad (26)$$

There is another big source of T parameter contribution from the heavy Higgs [8], with the reference higgs mass of $m_{h_{ref}} = 120$ GeV:

$$\Delta T_h = -\frac{3}{16\pi c^2} \log \left(\frac{m_h^2}{m_{h_{ref}}^2} \right). \quad (27)$$

Since in the interested region of parameter space this model gives a negligible contribution to the U parameter, we can fix $U=0$, therefore the experimental constraints for the T parameter is [9]:

$$T|_{U=0} = 0.10 \pm 0.08. \quad (28)$$

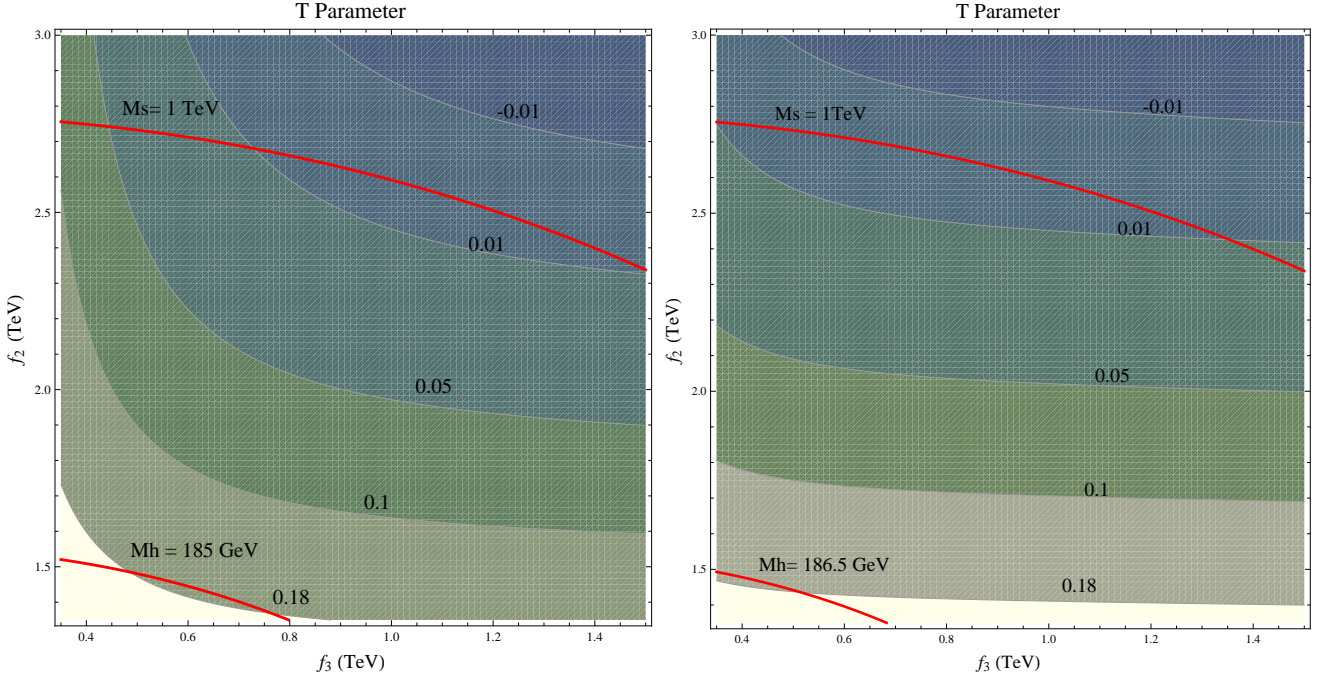


FIG. 2: Contours of T Parameter as functions of two VEVs f_2 and f_3 . Input parameters are : $m_{\tilde{t}_R} = 300\text{GeV}$, $\tan\beta = 2.0$, $\mu_2 = 5\text{TeV}$, $\mu_3 = 2\text{TeV}$, and $M_A = 800\text{GeV}$. In the left panel, the gauge couplings are $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$. In the right panel, the gauge couplings are $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.37$, $\hat{g}_{1V} = 2.5$, and $\hat{g}_2 = 0.78$. The bottom red lines in both plots put a lower bound for the Higgs bosons mass under the requirement of $T < 0.18$. The SUSY breaking scale $M_S = \sqrt{m_{\tilde{Q}}m_{\tilde{t}_R}}$ is required to be less than 1 TeV, corresponding to the top red lines in both plots, which put an upper bound for the parameter space.

The presence of a heavy Higgs boson with a mass much larger than 120 GeV will give a negative contribution to T parameter which may lead to a conflict with the experimental constraints. It is good for us that the mixing of gauge bosons instead drives the T parameter in the positive direction so that two effects may balance with each other and we can go back into the consistent region in the S - T plane. In Fig. [2], the contribution to T parameter

$m_{\tilde{t}_R}$ (GeV)	300	350	400	450	500	550
m_h (GeV)	185	188	190.6	193	195	196.6

TABLE III: Input parameters are : $m_{\tilde{t}_R} = 300$ GeV, $\tan\beta = 2.0$, $\mu_2 = 5$ TeV, $\mu_3 = 2$ TeV, and $M_A = 800$ GeV. Gauge couplings are taken to be $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$ and $\hat{g}_2 = 0.78$. Increasing the right handed stop mass from 300 GeV to 550 GeV in a modest way without inducing large fine tuning, the respective lower bound for the Higgs bosons mass can be calculated by requiring $T < 0.18$. As we can see, m_h varies from 185GeV to 196.6GeV.

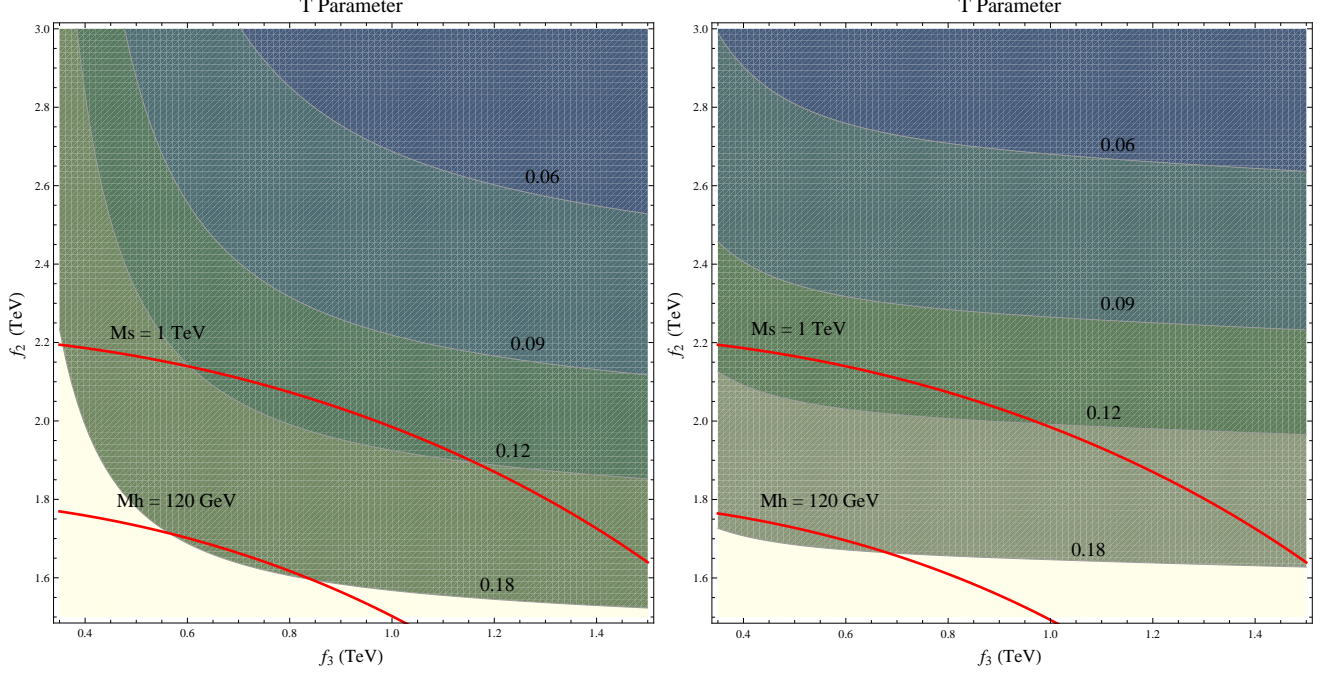


FIG. 3: Contours of T parameter as functions of two VEVs f_2 and f_3 . Input parameters are : $m_{\tilde{t}_R} = 300$ GeV, $\tan\beta = 0.9$, $\mu_2 = 5$ TeV, $\mu_3 = 2$ TeV, and $M_A = 400$ GeV. In the left panel, the gauge couplings are $\hat{g}_5 = 1.5$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$. In the right panel, the gauge couplings are $\hat{g}_5 = 1.5$, $\hat{g}_{1H} = 0.37$, $\hat{g}_{1V} = 2.5$, and $\hat{g}_2 = 0.78$. The bottom red lines in both plots put a lower bound for the Higgs bosons mass under the requirement of $T < 0.18$. The SUSY breaking scale $M_S = \sqrt{m_{\tilde{Q}}m_{\tilde{t}_R}}$ is required to be less than 1 TeV, corresponding to the top red lines in both plots, which put an upper bound for the parameter space.

combining both Eq. [26] and Eq. [27] is plotted against the VEVs of f_2 and f_3 . The lowest contour corresponds to the positive $T = 0.18$ bound. The requirement of $T < 0.18$ gives a lower bound to both the mass of vector top partner and the mass of the Higgs bosons. Since we demand less percentage of fine tuning, the SUSY breaking scale defined as $M_S = \sqrt{m_{\tilde{Q}}m_{\tilde{t}_R}}$ is required to be less than 1 TeV. This latter requirement fixes the upper bound for Higgs bosons in this model. In the left panel of Fig. [2], the gauge coupling constants are taken to be $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$. The mass of right handed stop is 300 GeV, μ -term masses are $\mu_2 = 5$ TeV, $\mu_3 = 2$ TeV and the input mass parameter for Higgs bosons in Eq. [20] is $M_A = 800$ GeV. The $\tan\beta$ is taken to be 2.0. This set of chosen parameters predicts the Higgs mass to be in the range of (185 GeV, 198 GeV), related to a vector top partner with its mass in the range of (1.85 TeV, 3.3 TeV). In the right panel of Fig. [2], we take another set of gauge coupling constants $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.37$, $\hat{g}_{1V} = 2.5$, and $\hat{g}_2 = 0.78$, with the other input parameters unchanged. As we can see, increasing \hat{g}_{1V} gauge coupling will reduce the T parameter's dependence on the value of f_3 parameter, i.e. the contour becomes more flat in the right panel, which will relax the lower bound of the vector top partner to be $m_{\tilde{Q}} > 1.82$ TeV and therefore result in less radiative correction to the Higgs mass. However increasing \hat{g}_{1V}

gauge coupling at the same time enlarges the tree level Higgs quartic coupling through the Δ_Y parameter so that the total effect is that with a larger \hat{g}_{1V} coupling the Higgs bosons mass is increased by just 1 – 2 GeV to be located in a range of (186.5 GeV, 200 GeV) . Varying the mass of the right handed stop by hundred GeVs can result in the lower bound of the Higgs mass slightly changing by a few GeVs. As shown in Table III, for a specific set of the gauge couplings: $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$, when the right handed stop mass is varied from 300GeV to 550 GeV , the lower bound for the Higgs bosons mass which satisfies the requirement of $T < 0.18$ will change accordingly from 185GeV to 196.6GeV. The parameter space constrained by the oblique parameter does not exclude a light Higgs boson, which is preferred by current ATLAS and CMS search results. Recent LHC experiments observed an excess of events at 125GeV in the final state of $\gamma\gamma$ hence give a hint that a Standard Model like Higgs may exist in the mass window of 123GeV – 127GeV. The light Higgs scenario can be achieved by tuning the $\tan\beta$. When we take $\tan\beta = 0.9$, the Higgs mass is limited to be $m_h \geq 120\text{GeV}$ depending on the specific gauge couplings and other input parameters as shown in Fig. [3]. But it will require a heavier vector top partner with its mass larger than 2.7 TeV to be consistent with the T parameter constraints.

Another important constraint comes from the corrections to the $Z \rightarrow b\bar{b}$ vertex. The b_L quark in this model is a linear combination of several fields and is mostly the gaugino of $SU(5)$, while the right handed bottom quark residing in the first moose site is only gauged with $SU(3) \times SU(2) \times U(1)$. The left handed bottom quark couples to the heavy gauge bosons in a way different from the right handed bottom quark. The gauge bosons mixing through Higgs VEVs induces both a correction to $Zb_L\bar{b}_L$ coupling and a correction to $Zb_R\bar{b}_R$ coupling:

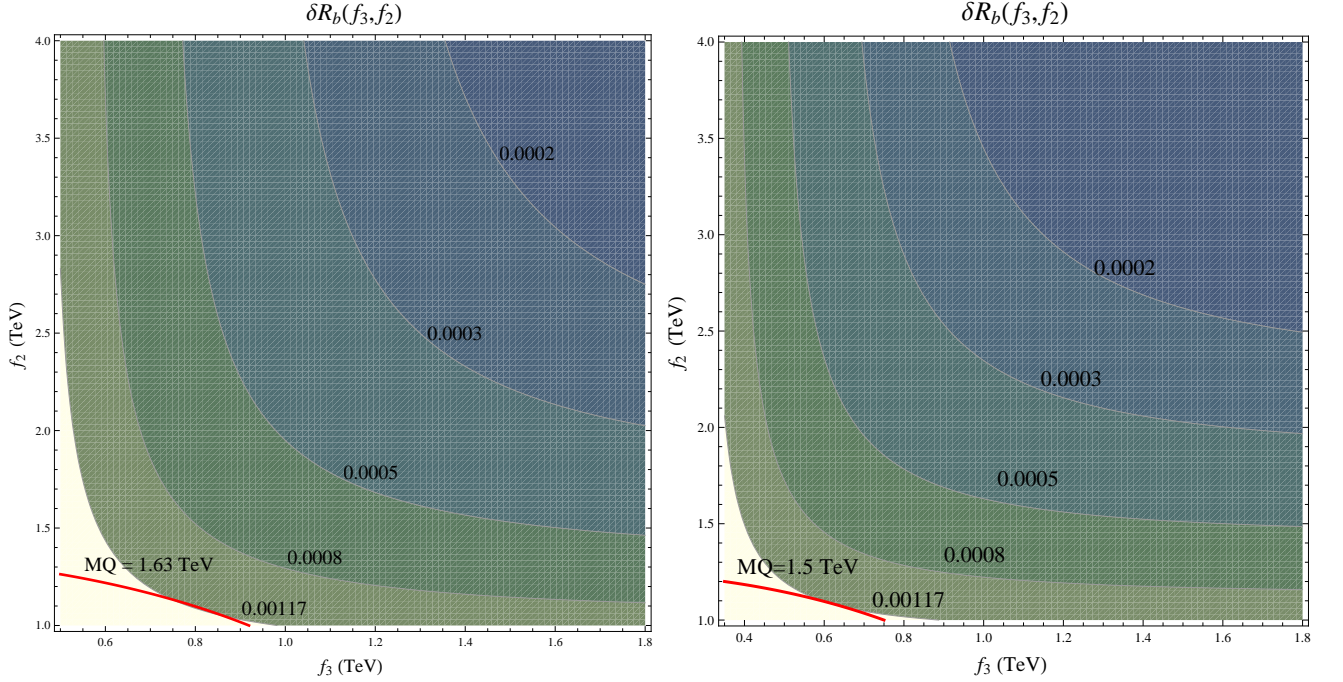


FIG. 4: contours of R_b as functions of two VEVs f_2 and f_3 . In the left panel, the gauge couplings are $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$. in the right panel the gauge couplings are $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.37$, $\hat{g}_{1V} = 2.5$ and $\hat{g}_2 = 0.78$. The red lines give the lower bound for the mass of the vector top partner under the requirement of $R_b < 0.00117$.

$$\begin{aligned}
\delta g_{zb_L b_L} = & - \frac{e}{sc} z_\mu J_3^\mu \left(\frac{\hat{g}_5^4}{(\hat{g}_2^2 + \hat{g}_5^2)^2} \frac{v^2}{8f_2^2} \right) - \frac{e}{sc} z_\mu J_Y^\mu \left(\frac{45\hat{g}_5^2 \hat{g}_{1H}^2 \hat{g}_{1V}^2 (15\hat{g}_5^2 + \hat{g}_{1V}^2)}{(15\hat{g}_5^2 \hat{g}_{1V}^2 + \hat{g}_{1H}^2 (15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{4f_3^2} \right) \\
& - \frac{e}{sc} z_\mu J_Y^\mu \left(\frac{15\hat{g}_5^2 (5\hat{g}_{1H}^2 + 3\hat{g}_{1V}^2) (15\hat{g}_5^2 \hat{g}_V^2 + \hat{g}_{1H}^2 (-15\hat{g}_5^2 + 4\hat{g}_{1V}^2))}{(15\hat{g}_5^2 \hat{g}_{1V}^2 + \hat{g}_{1H}^2 (15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{8f_2^2} \right) , \quad (29)
\end{aligned}$$

$$\begin{aligned} \delta g_{zb_R b_R} = & -\frac{e}{sc} z_\mu J_Y^\mu \left(\frac{3\hat{g}_{1H}^2 (10\hat{g}_5^2 - \hat{g}_{1V}^2) (15\hat{g}_5^2 \hat{g}_{1V}^2 + \hat{g}_{1H}^2 (-15\hat{g}_5^2 + 4\hat{g}_{1V}^2))}{(15\hat{g}_5^2 \hat{g}_{1V}^2 + \hat{g}_{1H}^2 (15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{8f_2^2} \right) \\ & + \frac{e}{sc} z_\mu J_Y^\mu \left(\frac{3\hat{g}_{1H}^4 (15\hat{g}_5^2 + \hat{g}_{1V}^2)^2}{(15\hat{g}_5^2 \hat{g}_{1V}^2 + \hat{g}_{1H}^2 (15\hat{g}_5^2 + \hat{g}_{1V}^2))^2} \frac{v^2}{4f_3^2} \right). \end{aligned} \quad (30)$$

The expression shows that the correction to Z gauge bosons coupled to the right handed bottom quark is much less as it is proportional to \hat{g}_{1H}^2 , and the value of \hat{g}_{1H} is assumed to be small in this model. Constraint from $Z \rightarrow b\bar{b}$ is measured by the branch ratio of $R_b = \Gamma(Z \rightarrow b\bar{b})/\Gamma(\text{hadron})$. The deviation of R_b due to the new physics can be expressed in terms of δg_L^{NP} and δg_R^{NP} :

$$\delta R_b = 2R_b(1 - R_b) \left(\frac{g_L}{g_L^2 + g_R^2} \delta g_L^{NP} + \frac{g_R}{g_L^2 + g_R^2} \delta g_R^{NP} \right) \quad (31)$$

$$g_L = -\frac{1}{2} + \frac{1}{3}s^2 \quad g_R = \frac{1}{3}s^2 \quad (32)$$

here R_b is the SM value predicted by the electroweak fit and its value is $R_b = 0.21578 + 0.0005(-0.0008)$. The deviation δR_b , used to describe the difference between its observation value and the SM fit result, is given by the experimental measurement [9]:

$$\delta R_b = 0.00051 \pm 0.00066 \quad (33)$$

Substituting Eq. [29] and Eq. [30] into Eq. [31], and we plot the dependence of δR_b on the two VEV parameters f_2 and f_3 in Fig. [4]. The lowest contour in that figure corresponds to the upper bound of $\delta R_b = 0.00117$, which gives a loose bound on the mass of the vector top partner compared with the T parameter constraint. For comparison reason, we will adopt the same two sets of gauge couplings to evaluate the value of δR_b , as we do in analyzing the T parameter. In the left panel of Fig. [4], the gauge couplings are: $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.378$, $\hat{g}_{1V} = 1.5$, and $\hat{g}_2 = 0.78$, by requiring $\delta R_b < 0.00117$, the mass of the vector top partner is limited to be $m_{\tilde{Q}} \geq 1.63$ TeV. In the right panel, gauge couplings are taken to be $\hat{g}_5 = 1.2$, $\hat{g}_{1H} = 0.37$, $\hat{g}_{1V} = 2.5$ and $\hat{g}_2 = 0.78$. Since both $\delta g_{zb_L b_L}$ and $\delta g_{zb_R b_R}$ will decrease as we increase the value of \hat{g}_{1V} , we could have a smaller bound value $m_{\tilde{Q}} \geq 1.5$ TeV for the vector top partner.

It can be seen that the T parameter measuring the amount of custodial symmetry breaking constrains the parameter space in a more stringent way. The measurement from $Z \rightarrow b\bar{b}$ consistently gives a comparable lower bound for the mass of the vector top partner in this model. Assuming the \hat{g}_{1V} coupling is relatively big i.e. $\hat{g}_{1V} \sim (1.5 - 2.5)$, the vector top partner mass $m_{\tilde{Q}}$ is required to be at least larger than 1.8 TeV and the parameter space is capable to accommodate a Higgs bosons with its mass in the range of (120 GeV, 200 GeV), after considering the electroweak constraints.

IV. CONCLUSIONS

In this paper, I present that adding extra singlet chiral superfields to interact with the link fields can trigger the gauge symmetry breaking and with an appropriately arranged superpotential, the VEVs will lead to spontaneous Supersymmetry breaking at the same time. I also add two chiral superfields transforming under the $SU(2)$ adjoint representation and the $SU(3)$ adjoint representation respectively to lift the moduli so that there is no light mode after the gauge symmetry breaking. Due to the nondecoupling D-term effects of those heavy fields, a larger Higgs quartic coupling is obtained. We explicitly demonstrate that in the large SUSY breaking limit, the effective low energy D-term is the same as in the unbroken gauge theory. Since the gauge couplings of the extra gauge groups $SU(5) \times U(1)_{1V}$ under which the Higgs boson are charged are taken to be strong, with a moderate $O(1)$ $\tan\beta$ the Higgs mass is heavy enough at the tree level. After taking the radiative corrections into account, the Higgs mass can be raised to be well beyond the LEP bound.

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